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R-parity violating anomaly mediated supersymmetry breaking

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Abstract: We propose a new scenario that solves the slepton negative mass squared problem of the minimal supersymmetric standard model with anomaly mediated supersymmetry breaking. The solution is achieved by including three trilinear R-parity violating operators in the superpotential. The soft supersymmetry breaking terms satisfy renormalisation group invariant relations in terms of supersymmetric couplings and the overall supersymmetry breaking mass scale. Flavour changing neutral currents can be naturally highly suppressed. A specific model predicts $\tan\beta = 4.2 \pm 1.0$. Excluding sleptons, the supersymmetric particle spectrum then depends upon two remaining free parameters. In the case of the R-parity violating couplings set at their quasi-fixed points at a supersymmetric GUT scale, the whole sparticle spectrum approximately depends upon only one free parameter. Imposing experimental limits leads to a constrained and distinctive phenomenology. The lightest CP-even Higgs of mass $m_h = 118$ GeV would be seen at the Tevatron. All sparticles and heavy Higgs would evade detection except for the lightest charginos and neutralinos, whose distinctive leptonic decays would be seen at the LHC.

Keywords: Supersymmetry Breaking, Beyond Standard Model, Supersymmetric Models.

Contents

1. Introduction

Low energy supersymmetry remains the most promising known perturbative solution to the gauge hierarchy problem that afflicts the Standard Model. It is clear from current data however, that for supersymmetry to be present in nature, it must be broken. Three phenomenologically distinct mechanisms for translating supersymmetry breaking from a hidden sector to the observable sector are currently recognised: tree-level gravity, gauge or anomaly mediation. The last mediator has received relatively little attention, and it is upon this mechanism that we focus the attention of this letter.

Anomaly mediated supersymmetry breaking (AMSB) in the minimal supersymmetric standard model (MSSM) provides a potential solution to the supersymmetric (SUSY) flavour problem [1]. This is a problem of many supergravity theories in which squarks and sleptons typically acquire unacceptably large flavour-changing neutral currents (FCNCs) through flavour mixings in their mass matrices. In AMSB, SUSY breaking is assumed to take place in a hidden ("sequestered") sector. A re-scaling anomaly in the super-Weyl conformal transformation transmits the SUSY breaking to the observable sector. It was suggested that the MSSM superfields be confined to a 3-brane in a higher dimensional bulk space-time, with the sequestered sector separated by the extra dimension from the visible sector brane. If direct Kahler couplings between the sequestered and visible sectors are suppressed (as is the case in the above geometrical set-up), these SUSY breaking terms can be the dominant forms of SUSY breaking in the visible sector. This scenario produces a supersymmetric spectrum dependent upon only three unknown parameters, an overall supersymmetric breaking mass scale and the MSSM Higgs potential parameters μ and B. For example, the AMSB mass squared values for scalar components of chiral matter supermultiplets are given by [1],

$$(m^2)^{\Phi_j}_{\Phi_i}|_{AM} = -\frac{1}{4}\ M^2_{\text{aux}}\Big[\mu\frac{d^2}{d\mu^2}\ln Z^j_i\Big], \tag{1.1}$$

where μ denotes the t'Hooft renormalization scale and Z^i_j is the matter field wave function of the superfield Φ_i. M_{aux} is the vacuum expectation value of a compensator superfield [1], and sets the overall mass scale for visible sector SUSY breaking. Defining

$$\Gamma^{\Phi_j}_{\Phi_i} \equiv -1/2\frac{d(\ln Z^i_j)}{d\ln\mu}, \tag{1.2}$$

we may write

$$(m^2)^{\Phi_j}_{\Phi_i}|_{AM} = \frac{1}{2}\ M^2_{\text{aux}}\Big[\beta(Y)\frac{\partial}{\partial Y}\Gamma^{\Phi_j}_{\Phi_i} + \beta(g)\frac{\partial}{\partial g}\Gamma^{\Phi_j}_{\Phi_i}\Big] \tag{1.3}$$

summed over all Yukawa couplings Y and gauge couplings g. $\beta(x)$ represents the beta function $dx/d\ln\mu$ of coupling x. An interesting fact is that the AMSB soft terms

are valid to all orders in perturbation theory [2]. For scalars of the first two families, Yukawa couplings can be safely neglected and so the dominant terms in eq. (1.3) are those proportional to the gauge couplings. These, being family universal, highly suppress the most problematic FCNC processes and thus solve the SUSY flavour problem.

The trilinear soft term A_Y corresponding to Yukawa coupling Y is given by [3]

$$YA_Y = -\beta(Y)M_{\text{aux}}, \tag{1.4}$$

and the gaugino mass M_g associated with each gauge group of coupling g is [3]

$$gM_g = \beta(g)M_{\text{aux}}. \tag{1.5}$$

One particularly desirable feature is that eqs. (1.3-1.5) are renormalisation invariant [1], provided that there are no massive parameters present in the superpotential (for example from bilinear terms). This means that AMSB has the advantage over gravity and gauge mediation that the low energy phenomenology does not depend upon the renormalisation of parameters at high scales, where unknown corrections from new physics apply. Unfortunately in the minimal version of the AMSB MSSM, the sleptons have negative mass squared values, indicating that the true vacuum state of the model is not the desired electroweak one. There have been several successful attempts to fix this problem, for example positive bulk Standard Model singlet contributions to the scalar masses [1], non-decoupling effects [4], heavy vector-like messengers coupled to light modulus fields [5] have all been proposed. All of the above models have spoiled the desirable renormalisation group invariant feature, rendering the theories potentially sensitive to unknown ultra-violet effects.

In a recent paper [6], a model which solves the slepton problem in AMSB with SUSY breaking Fayet-Iliopoulos D-terms of an additional U(1) gauge symmetry was presented. Extra Standard-Model gauge singlet chiral superfields are added to the MSSM. The model has the advantage of soft terms that do not depend upon the ultra-violet physics [6],[7]. The model of ref. [8] extends the MSSM by 3 extra Higgs doublets, a vector-like pair of extra triplets and 4 new singlets near the weak scale. Large Yukawa couplings between the extra Higgs and MSSM leptons provide additional positive contributions to the slepton mass squareds in eq. (1.3), while preserving renormalisation group invariance of the mass relations. We note that the above attempts to solve the AMSB slepton problem have all had R-parity invariance as a feature.

Here, we make a new proposal which preserves the renormalisation invariance of the AMSB supersymmetry breaking mass relations and requires no superfields additional to those in the MSSM coupling directly to the visible sector. By considering a subset of trilinear R-parity violating ($\not{R}_p$) operators in the superpotential, we change the wave-function renormalisation of the sparticles, in particular providing new positive contributions to the slepton mass squared values. Throughout this work we

will assume the dogma of minimality with respect to solving the slepton problem in AMSB, for brevity and simplicity.

First, in section 2, we will classify models that simultaneously solve the MSSM AMSB tachyonic slepton problem while not generating dangerous lepton flavour violating operators. There emerges a scenario with 3 non-negligible lepton number violating ($\not{L}$)-operators only. In section 3, we then present the slepton masses in terms of the supersymmetric couplings explicitly and briefly discuss the other soft breaking terms. All other soft masses are equivalent to the R_p conserving AMSB MSSM to one-loop order. In section 4, we impose constraints upon the model, the most restrictive being from lepton non-universality. For the $\not{L}$-contributions to be sufficient to raise all slepton mass squared values above zero, we require some $\not{L}$-couplings of order 1. We demonstrate that this is not in conflict with current data if the scalar sparticles are rather heavy, above 1.2 TeV. We also present the sparticle spectra. In section 5, some implications for collider searches at the Tevatron and LHC are presented. Finally, we summarise the main features of the model, reviewing its successful features and noting possible future work in section 6.

2. Rescuing the AMSB MSSM with R-parity violation

We begin with the AMSB MSSM including general trilinear $\not{R}_p$-operators. We then identify the subset of operators which are useful in solving the AMSB slepton problem. In the notation of ref. [9], the general trilinear $\not{R}_p$-MSSM superpotential is written

$$\begin{aligned} W_3 = {} & (Y_E)_{ij} L_i H_1 \bar{E}_j + (Y_D)_{ij} Q_i H_1 \bar{D}_j + (Y_U)_{ij} Q_i H_2 \bar{U}_j + \\ & \frac{1}{2}\lambda_{ijk} L_i L_j \bar{E}_k + \lambda'_{ijk} L_i Q_j \bar{D}_k + \frac{1}{2}\lambda''_{ijk} \bar{U}_i \bar{D}_j \bar{D}_k, \end{aligned} \tag{2.1}$$

where we have suppressed gauge indices, and $i, j, k, \ldots = 1, 2, 3$ are family indices. The anomalous dimensions $\Gamma_{\Phi_i}^{\Phi_j}$ relevant for substitution into eq. (1.3) for the superpotential W_3 have been presented in ref. [10] to two-loop order. Their one-loop truncation is annexed here to Appendix A for ease of reference.

For example, substituting $\Gamma_{E_{Ri}}^{E_{Rj}}$ into eq. (1.3), we obtain the most problematic soft mass squared, that of the right handed sleptons:

$$16\pi^2 (m^2_{E_R})^j_i = \frac{1}{2} M^2_{\text{aux}} \left[2(Y_E^\dagger)_{jk} \beta(Y_E)_{ki} + \lambda_{kmi} \beta(\lambda_{kmj}) + (i \leftrightarrow j) - G_E \right], \tag{2.2}$$

where $G_E \equiv 396 g_1^4/(25 \times 16\pi^2)$. The first term on the right hand side is negligible for selectrons and smuons because it is proportional to the electron and muon mass respectively. In the R-parity conserving scenario where all $\lambda_{ijk} = 0$, the last (negative) term therefore forces the right handed selectrons and smuons to have negative mass squared values. If $\tan\beta > 40$, the positive contribution from $(Y_E)_{33}$

becomes non-negligible and the right handed stau mass squared may be raised above zero. However, we immediately see that a positive contribution may be obtained to $(m_{E_R})^k_k$ from $\lambda_{ijk} \neq 0$, and it is this possibility that we exploit[1]. So far, the condition of minimality identifies[2] the combination

$$\lambda_{jk1}, \lambda_{lm2}, \lambda_{nq3} \neq 0, \tag{2.3}$$

which provides positive contributions to all three right handed slepton masses. We also observe that λ'_{ijk}, λ''_{ijk} cannot solve the problem of negative right handed slepton masses and so we drop them from the discussion. We will assume in the present models that they are zero, and indeed this assumption will make it simpler to satisfy stringent empirical limits upon successful scenarios.

The left-handed sleptons also have negative mass squared values in the usual R_p-conserving AMSB MSSM scenario. Including LLE operators by substituting $\Gamma^{L_j}_{L_i}$ into eq. (1.3) and setting $\lambda'_{ijk}, \lambda''_{ijk} = 0$ for all i, j, k,

$$16\pi^2 (m^2_L)^j_i = \frac{1}{2} M^2_{\mathrm{aux}} \left[(Y^\dagger_E)_{jk} \beta(Y_E)_{ki} + \lambda_{ikq}\beta(\lambda_{jkq}) + (i \leftrightarrow j) - G_L \right], \tag{2.4}$$

where $G_L \equiv (99g_1^2 + 75g_2^2)/(25 \times 16\pi^2)$. Thus the R-parity conserving scenario (all $\lambda_{ijk} = 0$) suffers from negative mass squared values for the left-handed selectron and smuon (and the stau if $\tan\beta$ is not large), analogous to the right handed sleptons. A positive contribution to $(m^2_L)^i_i$ results if $\lambda_{ijk} \neq 0$, i.e. we require

$$\lambda_{1jk}, \lambda_{2lm}, \lambda_{3nq} \neq 0, \tag{2.5}$$

to provide additional positive contributions to all left-handed slepton mass squareds.

In order to keep the number of couplings to a minimum, we require that the same non-zero couplings that render the right-handed slepton masses squared values positive in eq. (2.3) also provide us with positive left-handed slepton mass squared values in eq. (2.5).

We now identify a further constraint upon λ_{ijk}

$$\lambda_{imm} = 0, \quad \forall i \tag{2.6}$$

(no sum on m) to avoid the generation of large off-diagonal slepton mass terms. Such terms would generate an empirically unacceptable amount of lepton flavour violation [11], such as $\mu \to e\gamma$. Eq. (2.6) also forbids the generation of lepton-Higgs mixing, as can be seen from eq. (A.9). Simultaneously applying the above constraints in eqs. (2.3-2.6) leads to the unique combination

$$\lambda_{123}, \lambda_{132}, \lambda_{231} \neq 0. \tag{2.7}$$

[1] Following the minimality dogma, we assume real $\not{R}_p$-couplings.

[2] We set all $\not{R}_p$ operators not explicitly mentioned to zero.

Thus we have identified the $\not{L}$-couplings that will solve the AMSB MSSM slepton problem without generating lepton flavour violating effects that are too large. We note that for high $\tan\beta > 40$, we could set $\lambda_{123} = 0$ and still have positive mass squared values for the sleptons. For the moment, we include all three couplings for generality, and indeed below, we focus on a particular model for which $\tan\beta = 4.2$, requiring us to include $\lambda_{123} \neq 0$.

3. Soft supersymmetry breaking terms

We now discuss the (one loop) equations for the soft supersymmetry breaking terms. We work in a basis where Yukawa couplings apart from $Y_\tau \equiv (Y_E)_{33}$, $Y_t \approx (Y_U)_{33}$, $Y_b \approx (Y_D)_{33}$ (the tau, top and bottom Yukawa couplings respectively) and $\not{R}_p$-couplings *not* discussed above are sub-leading and are therefore neglected.

3.1 Slepton masses

The slepton soft masses are given by eqs. (2.4),(2.2) as

$$(m^2)^{L_1}_{L_1} = \frac{M^2_{aux}}{(16\pi^2)}\Big[\lambda_{123}\beta(\lambda_{123}) + \lambda_{132}\beta(\lambda_{132}) - \Big(\frac{3}{10}g_1\beta(g_1) + \frac{3}{2}g_2\beta(g_2)\Big)\Big] \tag{3.1}$$

$$(m^2)^{L_2}_{L_2} = \frac{M^2_{aux}}{(16\pi^2)}\Big[\lambda_{231}\beta(\lambda_{231}) + \lambda_{123}\beta(\lambda_{123}) - \Big(\frac{3}{10}g_1\beta(g_1) + \frac{3}{2}g_2\beta(g_2)\Big)\Big] \tag{3.2}$$

$$(m^2)^{L_3}_{L_3} = \frac{M^2_{aux}}{(16\pi^2)}\Big[Y_\tau\beta(Y_\tau) + \lambda_{132}\beta(\lambda_{132}) + \lambda_{231}\beta(\lambda_{231}) - \Big(\frac{3}{10}g_1\beta(g_1) + \frac{3}{2}g_2\beta(g_2)\Big)\Big]$$

$$(m^2)^{E_1}_{E_1} = \frac{M^2_{aux}}{(16\pi^2)}\Big[2\lambda_{231}\beta(\lambda_{231}) - \frac{6}{5}g_1\beta(g_1)\Big] \tag{3.3}$$

$$(m^2)^{E_2}_{E_2} = \frac{M^2_{aux}}{(16\pi^2)}\Big[2\lambda_{132}\beta(\lambda_{132}) - \frac{6}{5}g_1\beta(g_1)\Big] \tag{3.4}$$

$$(m^2)^{E_3}_{E_3} = \frac{M^2_{aux}}{(16\pi^2)}\Big[2Y_\tau\beta(Y_\tau) + 2\lambda_{123}\beta(\lambda_{123}) - \frac{6}{5}g_1\beta(g_1)\Big] \tag{3.5}$$

where the β-functions are given by

$$\beta(Y_\tau) = \frac{Y_\tau}{16\pi^2}\Big[4Y_\tau^2 + 3Y_b^2 + 2\lambda_{123}^2 + \lambda_{132}^2 + \lambda_{231}^2 - \Big(\frac{9}{5}g_1^2 + 3g_2^2\Big)\Big] \tag{3.6}$$

$$\beta(\lambda_{123}) = \frac{\lambda_{123}}{16\pi^2}\Big[2Y_\tau^2 + 4\lambda_{123}^2 + \lambda_{231}^2 + \lambda_{132}^2 - \Big(\frac{9}{5}g_1^2 + 3g_2^2\Big)\Big] \tag{3.7}$$

$$\beta(\lambda_{231}) = \frac{\lambda_{231}}{16\pi^2}\Big[Y_\tau^2 + 4\lambda_{231}^2 + \lambda_{123}^2 + \lambda_{132}^2 - \Big(\frac{9}{5}g_1^2 + 3g_2^2\Big)\Big] \tag{3.8}$$

$$\beta(\lambda_{132}) = \frac{\lambda_{132}}{16\pi^2}\Big[Y_\tau^2 + 4\lambda_{132}^2 + \lambda_{123}^2 + \lambda_{231}^2 - \Big(\frac{9}{5}g_1^2 + 3g_2^2\Big)\Big] \tag{3.9}$$

3.2 Other soft terms

The soft terms for squark masses and trilinear couplings can be derived from the general formulae in the Appendix. To one-loop accuracy, the rest are equivalent to the R_p conserving MSSM soft terms [12], except for the trilinear slepton couplings and m_{H_1}:

$$m_{H_1}^2 = \frac{M_{aux}^2}{16\pi^2}\Big[3Y_b\beta(Y_b) + Y_\tau\beta(Y_\tau) - \Big(\frac{3}{10}g_1\beta(g_1) + \frac{3}{2}g_2\beta(g_2)\Big)\Big]. \tag{3.10}$$

From eq. (3.6), we see that m_{H_1} depends upon the combination $\kappa \equiv 2\lambda_{123}^2 + \lambda_{132}^2 + \lambda_{231}^2$. Because μ is fixed partly by m_{H_1} in the electroweak symmetry breaking condition [13], it is altered from the R_p-conserving scenario by $\kappa \neq 0$. We note in particular the anomaly-mediated contribution to the B-term realised in a specific model with a bulk contribution [1]:

$$B = -\frac{M_{aux}}{16\pi^2}\Big[3Y_t^2 + 3Y_b^2 + Y_\tau^2 - \Big(\frac{3}{5}g_1^2 + 3g_2^2\Big)\Big]. \tag{3.11}$$

We shall utilise this model in order to cut the parameter space down. The prediction of B (for a given value of $M_{\rm aux}$) results in a prediction of $\tan\beta$ from the potential minimisation conditions [13]. We note that in the AMSB MSSM, a term $\mu H_1 H_2$ in the superpotential spoils the conformal invariance. However, μ can be viewed as a result of supersymmetry breaking [1], providing a natural explanation the size of μ necessary to obtain $M_Z = 91.18$ GeV. In our convention, B and μ have opposite signs in successful minima, so the B term predicted also constrains the sign of μ to be positive.

4. Spectra and constraints

Some products of the $\not{L}$-couplings are constrained to be tiny and practically useless in solving the AMSB slepton mass problem [9],

$$\lambda_{mni} \mapsto\!\!\!\!\!/\;\; \lambda_{mnj} \qquad i \neq j, \tag{4.1}$$

$$\lambda_{imn} \mapsto\!\!\!\!\!/\;\; \lambda_{jmn} \qquad i \neq j, \tag{4.2}$$

where $\mapsto\!\!\!\!\!/$ stands for 'not non-zero with'. The combination of couplings $\lambda_{123}, \lambda_{132}, \lambda_{231} \neq 0$ respects this constraint. In addition, the most recent bounds upon the individual couplings are [10]:

$$\lambda_{123} \lesssim 0.49 \times \frac{m_{\tilde{\tau}_R}}{1\ \mathrm{TeV}} \tag{4.3}$$

$$\lambda_{132} \lesssim 0.62 \times \frac{m_{\tilde{\mu}_R}}{1\ \mathrm{TeV}} \tag{4.4}$$

$$\lambda_{231} \lesssim 0.70 \times \frac{m_{\tilde{e}_R}}{1\ \mathrm{TeV}}, \tag{4.5}$$

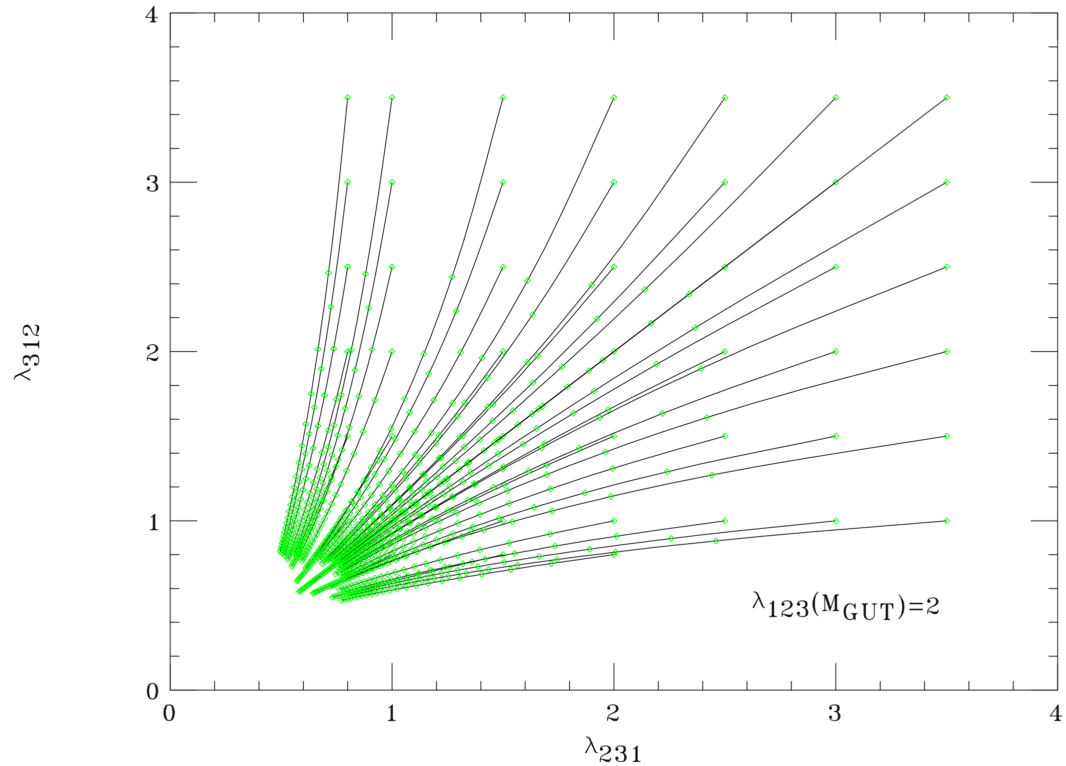

Figure 1: Quasi-fixed structure of the $\not R_p$-couplings. Renormalisation group trajectories of the couplings $\lambda_{132}, \lambda_{231}$ are shown for various boundary conditions at M_{GUT}. We display the running from M_{GUT} to the electroweak scale, where all couplings approach the $\lambda_{132} \sim 0.75$, $\lambda_{231} \sim 0.75$ region. The dots denote a decrease in the renormalisation scale by a factor of 100.

from charged lepton universality [14]. As we show below, these provide the most severe constraints upon the model. The couplings also pass the $\mu \to e\gamma$ conversion limits [15].

4.1 Sparticle Spectra

We now perform a one-loop accuracy numerical analysis to determine the sparticle and Higgs spectrum. The full one-loop Higgs potential is minimised at a scale $Q \equiv$2 TeV, where the radiative corrections to the potential are small, to determine μ. This choice of scale can change the $\tan\beta$ prediction a little. For experimental inputs on the gauge couplings we use [16] $\alpha_s(M_Z) = 0.119$, $\sin^2\theta_w = 0.2312$, $\alpha(M_Z) = 1/127.9$ in the $\overline{MS}$ scheme. 3 loop QCD$\otimes$2 loop QED is used as an effective field theory below M_Z. Central values [16] of the top and bottom pole masses $M_t = 174.3, M_b = 4.9$ GeV are taken. For simplicity we set $\lambda \equiv \lambda_{123} = \lambda_{231} = \lambda_{312}$, a renormalisation group invariant choice in the limit that $Y_\tau = 0$. In fact, all soft masses and couplings except those of the sleptons and one of the Higgs are independent of the choice of these three couplings. Here, we choose $\lambda(M_Z) = 0.73$, which is sufficient to render all

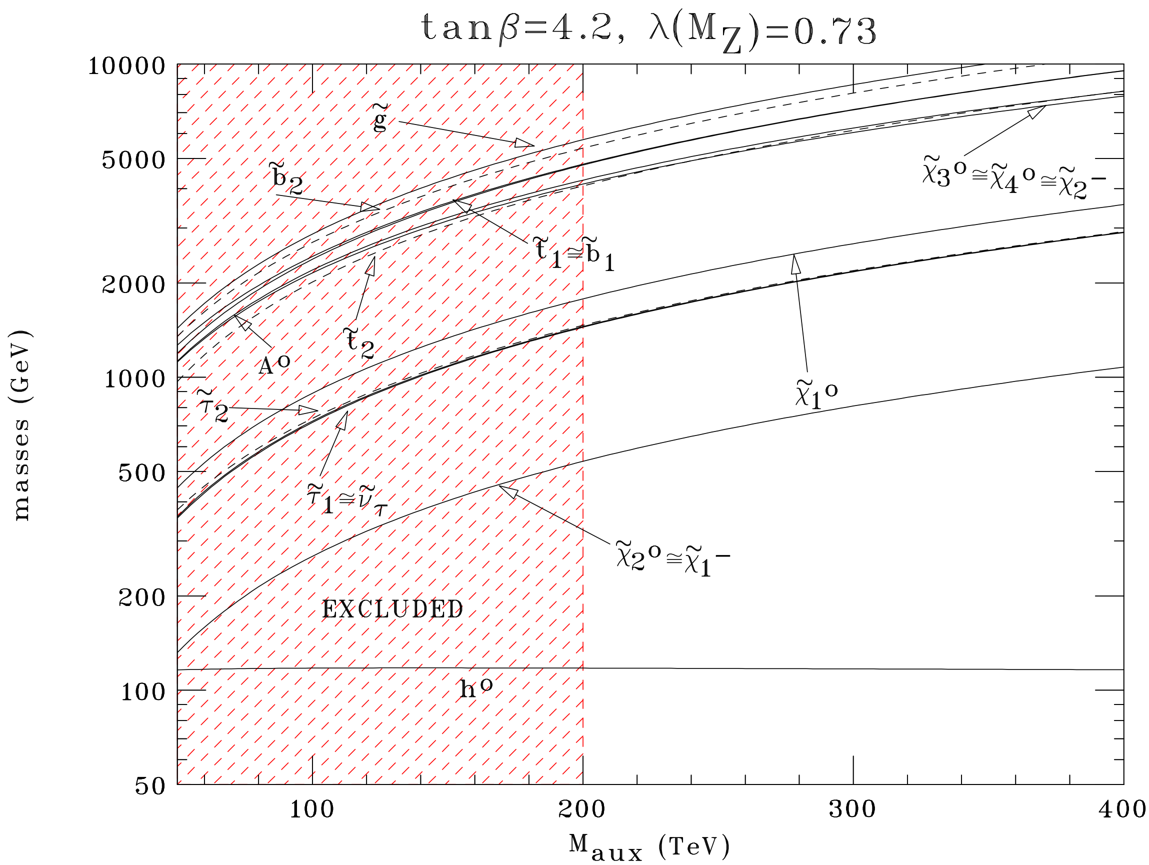

Figure 2: The supersymmetric particle spectrum in our AMSB scenario for different values of M_{aux}. The dashed region is excluded from charged lepton universality constraints.

slepton mass squared values positive ($\lambda > 0.66$). If $\lambda(M_Z)$ is set too large, then M_{aux} must be set very large in order to produce sleptons heavy enough to evade eqs. (4.3-4.5). Remarkably, $\lambda(M_Z) = 0.73$ is near a common quasi-fixed point value for all three couplings if we assume that they are set large at a scale $M_{GUT} \sim 2 \times 10^{16}$ GeV, where the gauge couplings unify. The quasi-fixed behaviour is exhibited by displaying insensitivity to the ultra-violet boundary condition [17]. This behaviour is exhibited in Fig. 1 As can be seen from the figure, λ_{132} and λ_{231} both approach the 0.7-0.8 region, roughly independent of the values assumed for them at M_{GUT}. We have checked that λ_{123} approaches the 0.7-0.8 region also.

Fig. 2 shows the supersymmetric particle spectrum in the AMSB scenario for different values of M_{aux}. The value of $\tan\beta \approx 4.2$ is predicted where the minimisation conditions of the Higgs potential satisfy eq. (3.11). In fact, this value has a small (neglected) dependence upon M_{aux} and the scale at which the potential is minimised, and thus has an uncertainty of about ± 1.0. The dashed region is excluded from the experimental bounds derived from charged lepton universality, the most stringent

$A_t(M_Z)$ =6.53	$A_b(M_Z)$ =11.97	$A_\tau(M_Z)$ =-0.71	$B(M_Z)$ =-1.42	$\mu(M_Z)$ =4.53
$m_{\tilde{g}}$ =6.31	$m_{\tilde{\chi}_1^0}$ =1.96	$m_{\tilde{\chi}_2^0}$ =0.59	$m_{\tilde{\chi}_3^0}$ =4.53	$m_{\tilde{\chi}_4^0}$ =4.53
m_h =0.1176	$m_A \approx m_H$ =4.67	$M_{H^\pm}$ =4.67	$m_{\tilde{\chi}_1^\pm}$ =0.59	$m_{\tilde{\chi}_2^\pm}$ =4.53
$m_{\tilde{\tau}_{1,2}}$ =1.60	$m_{\tilde{\nu}_\tau}$ =1.60	$m_{\tilde{e}_1,\tilde{\mu}_1}$ =1.60	$m_{\tilde{e}_2,\tilde{\mu}_2}$ =1.61	$m_{\tilde{\nu}_e,\tilde{\nu}_\mu}$ =1.60
$m_{\tilde{t}_1}$ =5.27	$m_{\tilde{t}_2}$ =4.50	$m_{\tilde{b}_1}$ =5.24	$m_{\tilde{b}_2}$ =5.93	
$m_{\tilde{u}_1,\tilde{c}_1}$ =5.86	$m_{\tilde{u}_2,\tilde{c}_2}$ =5.92	$m_{\tilde{d}_1,\tilde{s}_1}$ =5.86	$m_{\tilde{d}_2,\tilde{s}_2}$ =5.94	

Table 1: Spectrum and couplings for $M_{\rm aux}$ = 220 TeV and $\lambda(M_Z)$ = 0.73. Masses are measured in TeV.

being eq. (4.3). An improvement of these bounds by a factor of two would test up to M_{aux} = 400 TeV. The LSP neutralino is quasi-degenerate with the lightest chargino, as usual in AMSB [12]. The Higgs mass determinations are performed using state-of-the-art two loop corrections [18]. The lightest CP-even Higgs mass m_{h^0} is insensitive to $M_{\rm aux}$ and $\lambda(M_Z)$, but has the usual large dependence upon M_t. For $M_t = 174.3$ GeV however, $M_{h^0} = 117.5 \pm 0.5$ GeV, with an estimated 2 GeV uncertainty coming from higher order corrections. As can be seen from the figure, the excluded region forces all sparticles to be rather heavy - the lightest chargino and neutralino can be as light as 500 GeV, but all other sparticles and heavy Higgs must be heavier than 1.1 TeV.

5. Collider phenomenology

	σ(Pb)	No. events
	LHC	
$\chi_i^{\pm,0}$	0.025	720
$\tilde{q}, \tilde{g}$	10^{-11}	0
$\tilde{l}, \tilde{\nu}$	6×10^{-6}	0
	Tevatron	
All SUSY	2.3×10^{-6}	0

Table 2: Cross sections σ of sparticle production at the Tevatron and LHC.

The phenomenology and search prospects of the AMSB R_p-conserving scenarios have been considered in refs. [1],[19],[3],[12],[20]. The present scenario differs in two main respects. Firstly, the $\not{R}_p$-coupling exclusion limits shown in fig. 2 force superpartners to be heavier than was previously considered with R_p-conserving AMSB. Secondly, the decays of $\chi_1^\pm$, χ_1^0 are qualitatively different. In the R_p conserving case, the quasi-degeneracy of $\chi_1^\pm$ and χ_1^0 means that the $\chi_1^\pm$ is quasi-stable [19], and of course the χ_1^0 is undetected, except as missing energy. A classic signature for the lightest chargino is then the presence of heavily ionising tracks, with possible slow decays into pions/leptons. In the present scenario however, the lightest chargino and neutralinos decay almost immediately into 3 leptons.

To illustrate the decays and cross sections of the model, we pick a particular value of $M_{\rm aux}$ = 220 TeV. The detailed spectrum and parameters are displayed in Table 1. HERWIG6.1 [21] was utilised to estimate sparticle discovery prospects

of this spectrum at the Tevatron and LHC. We display the cross-sections of the hard sub-process of sparticle production in Table 2. Also shown is the number of expected events for luminosities of $\mathcal{L} = 10, 30$ fb^{-1} at the Tevatron and LHC respectively[3]. The table shows that charginos and neutralinos are produced at the LHC at a detectable rate, but the Tevatron should see no superparticles. However, we note that the lightest CP-even Higgs of mass $m_{h^0} \approx 118$ GeV should be discovered at the Tevatron [22].

We ran the weak-scale spectrum through a version of ISASUSY [23] modified to take $\not{R}_p$-interactions into account [24]. This then calculated the relevant decays. The lightest neutralino decays through twelve channels with equal branching ratios of 1/12 and partial widths of 2.25×10^{-5} GeV:

$$\begin{aligned} \chi_1^0 \to\ & e^+\bar{\nu}_\mu\tau^-,\ e^-\nu_\mu\tau^+,\ e^+\bar{\nu}_\tau\mu^-,\ e^-\nu_\tau\mu^+,\ \mu^+\bar{\nu}_e\tau^-,\ \mu^-\nu_e\tau^+, \\ & \mu^+\bar{\nu}_\tau e^-,\ \mu^-\nu_\tau e^+,\ \tau^+\bar{\nu}_e\mu^-,\ \tau^-\nu_e\mu^+,\ \tau^+\bar{\nu}_\mu e^-,\ \tau^-\nu_\mu e^+\ , \end{aligned} \tag{5.1}$$

whereas χ_1^+ has six decay channels

$$\chi_1^+ \to \nu_\mu\nu_e\tau^+,\ \nu_\tau\nu_e\mu^+,\ \nu_\tau\nu_\mu e^+,\ \mu^+e^+\tau^-,\ \tau^+e^+\mu^-,\ \tau^+\mu^+e^- \tag{5.2}$$

again with equal branching ratios of 1/6 and partial widths of 4.5×10^{-5} GeV. These decays should be easy to find with low backgrounds at the LHC. Double chargino production can be found by decays into six charged leptons, or chargino/neutralino production via a five charged lepton channel, with distinctive flavour structure. Lepton flavour violation is usually explicit in the final state. We have obtained approximately equal branching ratios here mainly because we have assumed $\lambda_{123} = \lambda_{231} = \lambda_{132}$. In the case they are non-degenerate, this will change and the relative branching ratios into different final states will help measure the $\not{R}_p$-couplings.

6. Conclusions

We have proposed a new solution to the problem of slepton negative mass squared values in the AMSB MSSM. It involves including 3 $\not{L}$-operators in the superpotential which were previously assumed to be absent. This leads to positive mass squared values for all of the sleptons and renormalisation-group invariant relations between supersymmetry breaking terms and the measured supersymmetric couplings. This has the advantage of rendering the model insensitive to unknown ultra-violet effects. The μ problem has a natural solution [1], indeed the prediction of the $B-$term in a particular model results in a prediction for $\tan\beta$. All of the sparticle spectrum except for the slepton masses are then given in terms of two free parameters: $M_{\rm aux}$ and κ.

[3]Equivalent to approximately 3 years running at low luminosity at the LHC

We have therefore assumed the MSSM spectrum near the weak scale, and that the dominant source of supersymmetry breaking terms in the observable sector is from anomaly mediation. Experimental limits on the $\not L$-operators provide stringent constraints upon the model, meaning that sparticle masses must be rather high. Whereas the lightest charginos and neutralinos can be as light as 500 GeV, the other sparticles must be heavier than 1.1 TeV. The Tevatron therefore sees no new particles except the lightest Higgs of mass 118 GeV, and the LHC can detect the lightest charginos and neutralinos via distinctive leptonic decays. Charged lepton universality violation is predicted to be close to its experimental bound, within a factor of two.

Neutrino masses and mixings are beyond the scope of this paper, but it is well known that the $\not L$-operators we have introduced can generate them at the loop level [25]. We intend to pursue them in future work [26], and the small number of free parameters should allow a strict correlation with lepton flavour violating predictions. We believe the present model of supersymmetry breaking in the observable sector to warrant several future works. For example, a more accurate calculation of the spectrum and a determination of the LHC reach in parameter space would be useful. It would be desirable to find symmetries to ban or suppress the other $\not R_p$-couplings. Aside from these, the usual calculations in the MSSM (quark FCNCs, charge and colour breaking constraints etc) could be performed. It will be interesting to investigate the present idea in a more general framework, for example when $\tan\beta$ is large (the prediction of the lightest MSSM Higgs mass is likely to change from the one presented here) and there are only two LLE couplings, or where splittings between the $\not R_p$couplings occur. Relaxing the assumption that $\lambda'_{ijk} = 0$ might lead to the possible observation of a single slepton at the LHC via slepton-strahlung [27].

To summarise, our scenario is a predictive scheme of supersymmetry breaking, containing natural solutions to the μ problem and supersymmetric flavour problem. The spectrum depends upon only two parameters apart from the slepton masses. In the case that the $\not R_p$-couplings are at their quasi-fixed point values, the slepton masses approximately only depend upon these same two parameters. If one assumes a high-energy cut-off scale, such as the GUT scale for example, we note that the weak-scale values of the couplings are approximately predicted by the quasi-fixed point and there is only one free parameter on which the whole sparticle spectrum depends. $\tan\beta$ is predicted in a specific model and the soft masses are given by renormalisation group invariant relations with the measured SUSY couplings. The phenomenology is rather distinctive and should be easily disentangled from other possibilities at the LHC, after being tested at the Tevatron by the Higgs mass prediction. The present model is the only example of a model with both insensitivity of the soft terms to unknown ultra-violet physics and the MSSM spectrum near the weak scale, and as such is important to investigate further.

A. One-loop anomalous dimensions and beta functions in the $R\!\!\!/_p$-MSSM

$\Lambda_{U^i}, \Lambda_{D^i}, \Lambda_{E^i}$ were written in a matrix notation in [10] as

$$(\Lambda_{E^k})_{ij} = \lambda_{ijk}, \qquad (\Lambda_{D^k})_{ij} = \lambda'_{ijk}, \qquad (\Lambda_{U^k})_{ij} = \lambda''_{ijk} \tag{A.1}$$

and we adopt this notation for presenting results with general family structures. The one-loop anomalous dimensions of the $R\!\!\!/_p$-MSSM are [10]

$$16\pi^2\ \Gamma^{(1)L_j}_{L_i} = \left(Y_E Y_E^\dagger\right)_{ji} + (\Lambda_{E^q}\Lambda^\dagger_{E^q})_{ji} + 3(\Lambda_{D^q}\Lambda^\dagger_{D^q})_{ji} - \delta^j_i(\frac{3}{10}g_1^2 + \frac{3}{2}g_2^2), \tag{A.2}$$

$$16\pi^2\ \Gamma^{(1)E_j}_{E_i} = 2\left(Y_E^\dagger Y_E\right)_{ji} + \mathrm{Tr}(\Lambda_{E^j}\Lambda^\dagger_{E^i}) - \delta^j_i(\frac{6}{5}g_1^2), \tag{A.3}$$

$$\begin{aligned}16\pi^2\ \Gamma^{(1)Q_j}_{Q_i} = &\left(Y_D Y_D^\dagger\right)_{ji} + \left(Y_U Y_U^\dagger\right)_{ji} + (\Lambda^\dagger_{D^q}\Lambda_{D^q})_{ij} \\ &-\delta^j_i(\frac{1}{30}g_1^2 + \frac{3}{2}g_2^2 + \frac{8}{3}g_3^2),\end{aligned} \tag{A.4}$$

$$\begin{aligned}16\pi^2\ \Gamma^{(1)D_j}_{D_i} = &\,2\left(Y_D^\dagger Y_D\right)_{ij} + 2\mathrm{Tr}(\Lambda^\dagger_{D^i}\Lambda_{D^j}) + 2(\Lambda_{U^q}\Lambda^\dagger_{U^q})_{ji} \\ &-\delta^j_i(\frac{2}{15}g_1^2 + \frac{8}{3}g_3^2)),\end{aligned} \tag{A.5}$$

$$16\pi^2\ \Gamma^{(1)U_j}_{U_i} = 2\left(Y_U^\dagger Y_U\right)_{ij} + \mathrm{Tr}(\Lambda_{U^j}\Lambda^\dagger_{U^i}) - \delta^j_i(\frac{8}{15}g_1^2 + \frac{8}{3}g_3^2)), \tag{A.6}$$

$$16\pi^2\ \Gamma^{(1)H_1}_{H_1} = \mathrm{Tr}\left(3Y_D Y_D^\dagger + Y_E Y_E^\dagger\right) - (\frac{3}{10}g_1^2 + \frac{3}{2}g_2^2), \tag{A.7}$$

$$16\pi^2\ \Gamma^{(1)H_2}_{H_2} = 3\mathrm{Tr}\left(Y_U Y_U^\dagger\right) - (\frac{3}{10}g_1^2 + \frac{3}{2}g_2^2), \tag{A.8}$$

$$16\pi^2\ \Gamma^{(1)H_1}_{L_i} = 16\pi^2\ \Gamma^{(1)L_i\,*}_{H_1} = -3(\Lambda^*_{D^q}Y_D)_{iq} - (\Lambda^*_{E^q}Y_E)_{iq}. \tag{A.9}$$

The beta functions of the couplings appearing in the superpotential in eq. (2.1) are:

$$\beta(Y_E)_{ij} = (Y_E)_{ik}\Gamma^{E_j}_{E_k} + (Y_E)_{ij}\Gamma^{H_1}_{H_1} - (\Lambda_{E^j})_{ki}\Gamma^{H_1}_{L_k} + (Y_E)_{kj}\Gamma^{L_i}_{L_k}, \tag{A.10}$$

$$\beta(Y_D)_{ij} = (Y_D)_{ik}\Gamma^{D_j}_{D_k} + (Y_D)_{ij}\Gamma^{H_1}_{H_1} - (\Lambda_{D^j})_{ki}\Gamma^{H_1}_{L_k} + (Y_D)_{kj}\Gamma^{Q_i}_{Q_k}, \tag{A.11}$$

$$\beta(Y_U)_{ij} = (Y_U)_{ik}\Gamma^{U_j}_{U_k} + (Y_U)_{ij}\Gamma^{H_2}_{H_2} + (Y_U)_{kj}\Gamma^{Q_i}_{Q_k}, \tag{A.12}$$

$$\begin{aligned}\beta(\Lambda_{E^k})_{ij} = &(\Lambda_{E^l})_{ij}\Gamma^{E_k}_{E_l} + (\Lambda_{E^k})_{il}\Gamma^{L_j}_{L_l} + (Y_E)_{ik}\Gamma^{L_j}_{H_1} \\ &-(\Lambda_{E^k})_{jl}\Gamma^{L_i}_{L_l} - (Y_E)_{jk}\Gamma^{L_i}_{H_1},\end{aligned} \tag{A.13}$$

$$\beta(\Lambda_{D^k})_{ij} = (\Lambda_{D^l})_{ij}\Gamma^{D_k}_{D_l} + (\Lambda_{D^k})_{il}\Gamma^{Q_j}_{Q_l} + (\Lambda_{D^k})_{lj}\Gamma^{L_i}_{L_l} - (Y_D)_{jk}\Gamma^{L_i}_{H_1}, \tag{A.14}$$

$$\beta(\Lambda_{U^i})_{jk} = (\Lambda_{U^i})_{jl}\Gamma^{D_k}_{D_l} + (\Lambda_{U^i})_{lk}\Gamma^{D_j}_{D_l} + (\Lambda_{U^l})_{jk}\Gamma^{U_i}_{U_l}. \tag{A.15}$$

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